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Heavy meson effective theory with $1/M_Q$ correction

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Abstract

We construct an effective Lagrangian of heavy and light mesons with $1/M_Q$ correction. The Lagrangian is constructed model independent way by using only the information of the symmetry of QCD. Reparameterisation invariance at the meson level is taken into account for the consistency of the theory. The partial decay width of the process $D^{*+} \to D^0\pi^+$ and the form factors of the process $\bar{B}^0 \to \pi^+ l\bar{\nu}$ are calculated with the $1/M_Q$ correction. We also introduce the light vector mesons based on the approach of the hidden local symmetry. The form factors of the process $\bar{B}^0 \to \rho^+ l\bar{\nu}$ are calculated with thus introduced ρ-meson. These results are easily translated to the D-meson semileptonic decays.

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The spin-flavor symmetry in the heavy quark sector [1] has been extensively studied to understand the heavy hadron decays. Since the proposed projects (B-factories, τ-charm factories and so on) are mainly accessible for the heavy mesons, we concentrate on the physics of the heavy mesons in this letter. There are two ways to extract the consequences of the symmetry. One is based on the heavy quark effective fields (heavy quark effective theory) [2], and another is based on the heavy meson effective fields (heavy meson effective theory) [3]. The algebra of the currents and charges in the heavy quark effective theory is used to reduce the number of independent form factors in the weak current matrix elements. It is well known that the six form factors in $B \to D$ and $B \to D^*$ weak transition matrix element are described by a single function called Isgur-Wise function in the heavy mass limit [1]. The $1/M_Q$ correction and the QCD correction of the heavy-to-heavy weak current matrix element have been calculated [4]. The reparameterisation invariance [5] plays an important role in the calculation. The result contains some form factors as unknown functions.

The construction of the heavy meson effective Lagrangian is based on both chiral symmetry in the light quark (u,d, and s quarks) sector and the spin-flavor symmetry in the heavy quark (c and b quarks) sector [3]. The semileptonic decays of the heavy mesons have been treated in leading order of the $1/M_Q$ expansion by using the heavy meson effective Lagrangian [3,6,7]. It should be noted that there is a B^* pole contribution in the form factors of $B \to \pi l \nu$, which is similar to the ρ pole contribution in the form factor of π. The B^* meson is naturally introduced as the spin partner of the B meson. Light vector mesons are introduced in the heavy meson effective Lagrangian [8] based on the hidden local symmetry [9].

In this paper we construct the heavy meson effective Lagrangian up to the $O(1/M_Q^2)$ corrections in $1/M_Q$ expansion and $O(p^2)$ in the chiral expansion [1]. The explicit chiral sym-

[1] H-Y.Cheng et al. have also considered the $1/M_Q$ correction in heavy meson effective Lagrangian [10]. But they have not emunerated all possible $O(1/M_Q)$ terms in terms of the effective meson field H_v, and also have not considered the $1/M_Q$ correction in weak currents.

metry breaking is worth studying, but it is left for our future works. Since the construction is based only on the symmetry alone, the results from the Lagrangian should be model independent. The Lagrangian contains ten parameters (six couplings, two heavy quark masses, a mass scale of the order of Λ_{QCD} (typical scale of the QCD dynamics), and a parameter which describe the mass splitting between the heavy pseudoscalar and vector mesons) which should be fixed by the experiments. The decay width of the process $D^{*+} \to D^0\pi^+$ and the form factors of the process $\bar{B}^0 \to \pi^+ l\bar{\nu}$ are calculated with the $1/M_Q$ correction. The light vector mesons are introduced according to the method of the hidden local symmetry. The form factors of the process $\bar{B}^0 \to \rho^+ l\bar{\nu}$ are also calculated with thus introduced light vector mesons.

The Lagrangian is constructed by the following effective fields. The heavy mesons are described by the field

$$H_v = \frac{1+ \not{v}}{2} \left[i\gamma_5 P_v + \gamma_\mu P_v^{*\mu} \right], \tag{1}$$

where v is the velocity of the heavy quark inside. The meson momentum p is described by

$$p = M_Q v + k, \tag{2}$$

where k is the residual momentum of the order of Λ_{QCD}. The fields P_v and P_v^*, which have mass dimension 3/2, are the heavy pseudoscalar and heavy vector fields, respectively. They are the multiplets in quark flavor given as

$$P_v^{(*)} = \begin{pmatrix} D^0 & D^+ & D_s^+ \\ B^- & \bar{B}^0 & \bar{B}_s^0 \end{pmatrix}^{(*)} . \tag{3}$$

The field H_v is transformed under the spin-flavor $SU(4)$ transformation which is decomposed by $SU(2)_{spin}$ and $SU(2)_H$ transformation and chiral transformation ($SU(3)_L \times SU(3)_R$ transformation) as

$$H_v \to S H_v, \tag{4}$$

$$H_v \to z_H H_v, \tag{5}$$

$$H_v \to H_v h(\Pi, g_L, g_R)^\dagger, \tag{6}$$

where $S \in SU(2)_{spin}$ acts on the Dirac index, $z_H \in SU(2)_H$ on the heavy flavor index. The chiral transformation is non-linearly realized as

$$\xi \to g_L \xi h(\Pi, g_L, g_R)^\dagger = h(\Pi, g_L, g_R)\xi g_R, \tag{7}$$

where $g_L \in SU(3)_L$, $g_R \in SU(3)_R$, and ξ is defined in terms of the field of the Nambu-Goldstone bosons (π, K, η_8),

$$\xi = e^{i\Pi/f_\pi} \qquad \left(\Pi = \Pi^a \frac{\lambda^a}{2}\right). \tag{8}$$

The unitary matrix $h(\Pi, g_L, g_R)$ has a complicated form, but if we consider the vector transformation $g = g_L = g_R$, then $h(\Pi, g_L, g_R) = g$. For the convenience to consider the chiral expansion (derivative expansion), the followings are used to build the Lagrangian,

$$\alpha_\perp^\mu = \frac{i}{2}\left(\xi\partial^\mu\xi^\dagger - \xi^\dagger\partial^\mu\xi\right), \tag{9}$$

$$\alpha_\parallel^\mu = \frac{i}{2}\left(\xi\partial^\mu\xi^\dagger + \xi^\dagger\partial^\mu\xi\right). \tag{10}$$

These are transformed under the chiral transformation as

$$\alpha_\perp^\mu \to h\alpha_\perp^\mu h^\dagger, \tag{11}$$

$$\alpha_\parallel^\mu \to h\alpha_\parallel^\mu h^\dagger + hi\partial^\mu h^\dagger. \tag{12}$$

The parity transformation and charge conjugation are defined as follows. The effective fields are transformed under the parity transformation as

$$\mathcal{P}H_v(x)\mathcal{P}^\dagger = \gamma^0 H_{\bar{v}}(\bar{x})\gamma^0, \tag{13}$$

$$\mathcal{P}\alpha_\perp^\mu(x)\mathcal{P}^\dagger = -\alpha_{\perp\mu}(\bar{x}), \tag{14}$$

$$\mathcal{P}\alpha_\parallel^\mu(x)\mathcal{P}^\dagger = \alpha_{\parallel\mu}(\bar{x}), \tag{15}$$

where $\bar{x} = (x^0, -\boldsymbol{x})$ and $\bar{v} = (v^0, -\boldsymbol{v})$, and under the charge conjugation

$$\mathcal{C}H_v(x)\mathcal{C}^\dagger = C\left(\overline{H_v^{(-)}}(x)\right)^T C^\dagger, \tag{16}$$

$$\mathcal{C}\alpha_\perp^\mu(x)\mathcal{C}^\dagger = \left(\alpha_\perp^\mu(x)\right)^T, \tag{17}$$

$$\mathcal{C}\alpha_\parallel^\mu(x)\mathcal{C}^\dagger = -\left(\alpha_\parallel^\mu(x)\right)^T, \tag{18}$$

where $\mathcal{P}$ and $\mathcal{C}$ are the operators, $C = i\gamma^2\gamma^0$, and $H_v^{(-)}$ is the effective field of the heavy mesons which contain the negative energy component of the heavy quark. The field $H_v^{(-)}$ is transformed in the same way as H_v under the spin and flavor transformation.

Now the effective Lagrangian can be obtained by imposing the chiral symmetry and the invariance under the parity transformation and charge conjugation. Spin-flavor symmetry is imposed with the breaking terms of $O(1/M_Q)$ included. We generally write down all the possible terms and get

$$
\begin{aligned}
\mathcal{L} = & -\sum_v \mathrm{tr}\left\{\bar{H}_v v\cdot iDH_v\right\} - \sum_v \mathrm{tr}\left\{\bar{H}_v \frac{(iD)^2}{2M} H_v\right\} \\
& +\Lambda\sum_v \mathrm{tr}\left\{\bar{H}_v H_v\right\} + \kappa'\Lambda\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{\Lambda}{M} H_v\right\} + \kappa\Lambda\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{\Lambda}{M}\sigma_{\rho\sigma} H_v \sigma^{\rho\sigma}\right\} \\
& +r\sum_v \mathrm{tr}\left\{\bar{H}_v H_v v\cdot\hat{\alpha}_{\parallel}\right\} \\
& \quad +r\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{iD_\mu}{2M} H_v \alpha_{\parallel}^{\mu}\right\} + h.c. \\
& \quad +r_1\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{\Lambda}{M} H_v v\cdot\hat{\alpha}_{\parallel}\right\} + r_2\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{\Lambda}{M}\sigma^{\rho\sigma} H_v \sigma_{\rho\sigma} v\cdot\hat{\alpha}_{\parallel}\right\} \\
& +\lambda\sum_v \mathrm{tr}\left\{\bar{H}_v H_v \gamma_\mu\gamma_5\alpha_{\perp}^{\mu}\right\} \\
& \quad -\lambda\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{v\cdot iD}{2M} H_v \gamma_\mu\gamma_5\alpha_{\perp}^{\mu}\right\} + h.c. \\
& \quad -\lambda\sum_v \varepsilon^{\mu\nu\rho\sigma}\mathrm{tr}\left\{\bar{H}_v \frac{iD_\rho}{4M} H_v \sigma_{\mu\nu}\alpha_{\perp\sigma}\right\} + \mathrm{h.c.} \\
& \quad +\lambda_1\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{\Lambda}{M} H_v \gamma_\rho\gamma_5\alpha_{\perp}^{\rho}\right\} + \lambda_2\sum_v \mathrm{tr}\left\{\bar{H}_v \frac{\Lambda}{M}\gamma_\rho\gamma_5 H_v \alpha_{\perp}^{\rho}\right\} \\
& +(\text{Anti-particle}),
\end{aligned} \tag{19}
$$

where $1/M = \mathrm{diag}(1/M_c, 1/M_b)$. Anti-particle part has exactly the same form of the particle part but $H_v \to H_v^{(-)}$ and $v \to -v$. The pseudoscalar meson masses m_P and the vector meson masses m_V are expanded as

$$
\begin{aligned}
m_P^2 &= M_Q^2\left\{1 + 2\frac{\Lambda}{M_Q} + 2\kappa'\frac{\Lambda^2}{M_Q^2} + 12\kappa\frac{\Lambda^2}{M_Q^2}\right\}, \\
m_V^2 &= M_Q^2\left\{1 + 2\frac{\Lambda}{M_Q} + 2\kappa'\frac{\Lambda^2}{M_Q^2} - 4\kappa\frac{\Lambda^2}{M_Q^2}\right\}.
\end{aligned} \tag{20}
$$

It should be noted that the normalisation of the effective field is different from that of the

conventional boson fields by $\sqrt{M_Q}$. Light vector mesons are introduced by the method of the hidden local symmetry [9]. The covariant derivative in the Lagrangian is defined as

$$iD_\mu H_v = i\partial_\mu H_v - H_v g_V V_\mu, \tag{21}$$

where $V_\mu = V_\mu^a \lambda^a/2$ is the light vector meson field, g_V is the gauge coupling constant of the hidden local symmetry, and $\hat{\alpha}_\parallel^\mu \equiv \alpha_\parallel^\mu - g_V V^\mu$. We used the equation of motion to drop out the several terms as the higher order terms in $1/M_Q$ expansion and chiral expansion. We take the reparameterisation invariance into account up to $O(1/M_Q^2)$ at the meson level. The reparameterisation transformation of the effective field is

$$\begin{aligned} H_v \to & \Lambda(v, \frac{w+iD/M}{|w+iD/M|})\Lambda(\frac{w+iD/M}{|w+iD/M|}, w) H_w \Lambda(w, \frac{w+i\overleftarrow{D}/M}{|w+i\overleftarrow{D}/M|})\Lambda(\frac{w+i\overleftarrow{D}/M}{|w+i\overleftarrow{D}/M|}, v) e^{-iqx} \\ &= \left\{ H_w - \frac{1}{2M}[\not{q}, H_w] \right\} e^{-iqx} + O(1/M^2), \end{aligned} \tag{22}$$

where $w = v + q/M$ and

$$\Lambda(w,v) = \frac{1+\not{w}\,\not{v}}{\sqrt{2(1+w\cdot v)}}. \tag{23}$$

Other fields are not transformed.

Many terms in the effective Lagrangian are dropped out by the reparameterisation invariance as follows: For example, we can consider the term

$$\mathrm{tr}\left\{ \bar{H}_v \frac{iD^\rho}{M} H_v \gamma_\rho \gamma_5 v\cdot \alpha_\perp \right\}. \tag{24}$$

The reparameterisation invariant form of this term is

$$\mathrm{tr}\left\{ \bar{\tilde{H}}_v \mathcal{V}^\rho \mathcal{V}^\sigma \tilde{H}_v \gamma_\rho \gamma_5 \alpha_{\perp\sigma} \right\}, \tag{25}$$

where $\mathcal{V} = v + iD/M$ and

$$\begin{aligned} \tilde{H}_v &= \Lambda(\frac{v+iD/M}{|v+iD/M|}, v) H_v \Lambda(v, \frac{v+i\overleftarrow{D}/M}{|v+i\overleftarrow{D}/M|}) \\ &= H_v + \frac{1}{2M}\left\{ [\gamma_\mu, iD^\mu H_v] - 2v\cdot iDH_v \right\} + O(1/M^2). \end{aligned} \tag{26}$$

Since $\tilde{H}_v$ satisfies the relations

$$\not{\mathcal{V}}\tilde{H}_v = \tilde{H}_v, \tag{27}$$

$$\tilde{H}_v \overleftarrow{\not{\mathcal{V}}} = -\tilde{H}_v, \tag{28}$$

the term (25) vanishes. Another example is the term

$$\epsilon^{\alpha\beta\rho\sigma}\mathrm{tr}\left\{\bar{H}_v\frac{iD_\alpha}{M}\sigma_{\rho\sigma}H_v\alpha_{\perp\beta}\right\}. \tag{29}$$

The reparameterisation invariant form of this term is

$$\epsilon^{\alpha\beta\rho\sigma}\mathrm{tr}\left\{\bar{\tilde{H}}_v\mathcal{V}_\alpha\sigma_{\rho\sigma}\tilde{H}_v\alpha_{\perp\beta}\right\}. \tag{30}$$

The leading component

$$\epsilon^{\alpha\beta\rho\sigma}\mathrm{tr}\left\{\bar{H}_v v_\alpha\sigma_{\rho\sigma}H_v\alpha_{\perp\beta}\right\} \tag{31}$$

which should be of the order of $1/M_Q$, because it breaks the spin symmetry. So the original term (29) is $O(1/M_Q^2)$. (The leading $O(1/M_Q)$ component reduces to the λ_2 term.)

We can do in the same way for the weak currents. The heavy-to-light weak current is obtained up to the $O(1/M_Q^2)$ corrections in $1/M_Q$ expansion and $O(p^2)$ in chiral expansion as

$$\begin{aligned}
J^{ia}_\mu(0) = F\Big[&\mathrm{tr}\left\{(\xi^\dagger)^{ji}\gamma_\mu(1-\gamma_5)H^{aj}_v\right\}\\
&+\frac{1}{2M_a}\mathrm{tr}\left\{(\xi^\dagger)^{ji}\gamma_\mu(1-\gamma_5)[\gamma_\rho, iD^\rho H^{aj}_v]\right\}\Big]\\
&+\alpha_1\frac{\Lambda}{M_a}\mathrm{tr}\left\{(\xi^\dagger)^{ji}\gamma_\mu(1-\gamma_5)H^{aj}_v\right\}\\
&+\alpha_2\frac{\Lambda}{M_a}\mathrm{tr}\left\{(\xi^\dagger)^{ji}\gamma_\mu(1-\gamma_5)\gamma^\rho H^{aj}_v\gamma_\rho\right\}\\
&+\beta_1\Lambda^{1/2}\mathrm{tr}\left\{(\xi^\dagger)^{ji}\gamma_\mu(1-\gamma_5)H^{aj}_v\left(v\cdot\hat{\alpha}_\parallel\right)\right\}\\
&+\beta_2\Lambda^{1/2}\mathrm{tr}\left\{(\xi^\dagger)^{ji}\gamma_\mu(1-\gamma_5)H^{aj}_v\left(\gamma_\rho\hat{\alpha}^\rho_\parallel\right)\right\}.
\end{aligned} \tag{32}$$

This current contains $L_\mu = \gamma_\mu(1-\gamma_5)$. We regard it in the construction of the current as an external field which is transformed under the spin-flavor transformation as

$$L_\mu \to g_L L_\mu z_H^\dagger, \tag{33}$$

$$L_\mu \to L_\mu S^\dagger. \tag{34}$$

We also impose the "parity" invariance under the transformation

$$L_\mu \to \gamma^0 L_\mu \gamma^0 \tag{35}$$

to keep the $V - A$ structure of the current.

From this current, we can extract the decay constants of the heavy mesons.

$$f_P = \sqrt{\frac{2}{M_Q}} \left\{ F + \frac{\Lambda}{M_Q}(\alpha_1 + 2\alpha_2) \right\}, \tag{36}$$

$$f_V = \sqrt{\frac{2}{M_Q}} \left\{ F + \frac{\Lambda}{M_Q}(\alpha_1 - 2\alpha_2) \right\}. \tag{37}$$

The heavy-to-heavy current can also be obtained in the same way, and the result contains some functions of $v \cdot v'$ instead of the parameters, where v and v' are the velocities of the heavy quarks. Weak interaction is introduced as the contact interaction between the currents [2].

We get the width of the $D^{*+} \to D^0\pi^+$ decay from the Lagrangian. Since the energy-momentum of π is very small (in the D^* rest frame), our Lagrangian is accessible for this process. We can use the axial vertex in the Lagrangian (couplings λ, λ_1, and λ_2). The width is calculated as

$$\Gamma(D^{*+} \to D^0\pi^+) = \frac{\lambda^2 M_c^2 \left(E_\pi^2 - m_\pi^2\right)^{3/2}}{12\pi m_{D^*}^2 f_\pi^2} \left[1 + \frac{E_\pi}{M_c} + \frac{2\left(\lambda_1 - \lambda_2\right)}{\lambda} \frac{\Lambda}{M_c} \right], \tag{38}$$

where

$$E_\pi = \frac{m_{D^*}^2 - m_D^2 + m_\pi^2}{2m_{D^*}}. \tag{39}$$

The term E_π/M_c can be neglected because it is $O(1/M_c^2)$ due to the kinematics ($E_\pi = O(\Lambda^2/M_c)$). But it should be kept in the off-shell amplitude as we will see next in $B \to \pi l \nu$

[2] Here, the QCD correction is not considered for the simplicity in the presentation of the Lagrangian and current construction.

decay. The couplings λ_1 and λ_2 are contained in the combination of $\lambda_1 - \lambda_2$. The branching ratio of the decay is already measured by CLEO [11]. If the total width is fixed, we can constraint the parameters λ and $\lambda_1 - \lambda_2$. (At present, we have only the upper bound by ACCMOR [12].)

The form factors of the process $\bar{B}^0 \to \pi^+ l \bar{\nu}$ can also be calculated. There is a B^*-pole contribution as in fig.1. The form factors are defined as

$$\langle \pi | J_\mu | \bar{B}^0 \rangle = f_+(q^2)(p_B + p_\pi)_\mu + f_-(q^2)(p_B - p_\pi)_\mu, \tag{40}$$

where $q^2 = (p_B - p_\pi)^2$. We obtain

$$f_\pm(q^2) = \frac{1}{2}\frac{f_B}{f_\pi}\left[1 - \frac{f_{B^*}}{f_B}\left\{\lambda\left(1 + \frac{v \cdot p_\pi}{2M_b}\right) + (\lambda_1 - \lambda_2)\frac{\Lambda}{M_b}\right\}\frac{2M_b\,(v \cdot p_\pi \mp m_B)}{q^2 - m_{B^*}^2}\right], \tag{41}$$

where

$$v \cdot p_\pi = \frac{m_B^2 + m_\pi^2 - q^2}{2m_B}. \tag{42}$$

The axial couplings λ_1 and λ_2 are contained in the combination of $\lambda_1 - \lambda_2$ also in this form factor. This result coincides with the result in ref. [13] in the soft pion limit $p_\pi \to 0$ [3]. Our result gives an extension of their result to $0 \leq |p_\pi| \lesssim 4\pi f_\pi$. (The upper bound $4\pi f_\pi$ is taken as the limit of the chiral expansion.)

This form factors are important in extracting $|V_{ub}|$ from the exclusive decay of $B \to \pi l \nu$. The q^2 dependence is given, which gives valuable information on the axial coupling constant by fitting the form factors with the q^2 spectrum.

The form factors of the process $\bar{B}^0 \to \rho^+ l \bar{\nu}$ can be calculated. There is a B-pole contribution as in fig.2. The form factors are defined as

$$\begin{aligned}\langle \rho | J^\mu | \bar{B}^0 \rangle &= f_1(q^2)\epsilon^{\mu\nu\rho\sigma}\epsilon^*_\nu(p_\rho)(p_B + p_\rho)_\rho(p_B - p_\rho)_\sigma \\ &+ f_2(q^2) i \epsilon^{*\mu}(p_\rho) \\ &+ f_3(q^2) i (\epsilon^*(p_\rho) \cdot p_B)(p_B + p_\rho)^\mu \\ &+ f_4(q^2) i (\epsilon^*(p_\rho) \cdot p_B)(p_B - p_\rho)^\mu,\end{aligned} \tag{43}$$

[3] The sign and the normalisation of the axial coupling λ is different from the one in ref. [13].

where $q^2 = (p_B - p_\rho)^2$. We obtain

$$f_1(q^2) = 0, \tag{44}$$

$$f_2(q^2) = f_B g_V - \beta_2 \sqrt{M_b \Lambda} g_V, \tag{45}$$

$$f_3(q^2) = \frac{1}{2}\beta_1 \frac{\sqrt{M_b\Lambda}}{M_b^2} g_V, \tag{46}$$

$$f_4(q^2) = 2 f_B g_V \left\{ r + \frac{\Lambda}{M_b}(r_1 + 6 r_2) \right\} \frac{1}{m_{B^\pm}^2 - q^2} + \frac{1}{2}\beta_1 \frac{\sqrt{M_b\Lambda}}{M_b^2} g_V. \tag{47}$$

As same as in the case of $\bar{B}^0 \to \pi^+ l \bar{\nu}$, this result is valid for the low energy ρ meson since we are using the chiral expansion. The form factor $f_4(q^2)$ is not significant as far as we can neglect the lepton masses. If we can fix the parameters β_1, β_2, f_B, and g_V, the Kobayashi-Maskawa matrix element V_{ub} can be extracted from this decay mode.

Many precise experiments which are expected in future B-factories will be used to fix the parameters in this effective theory. We have the predictive power once the parameters are fixed. The prediction is model independent, since we only used the information of the symmetry of QCD in constructing the effective theory.

In this paper, we construct the heavy meson effective Lagrangian up to $O(p^2)$ in the chiral expansion and $O(1/M_Q^2)$ in the $1/M_Q$ expansion. We enumerated all the possible terms which are allowed by the symmetry. The reparameterisation invariance in the meson level is very important in the construction. Many terms are dropped out by the invariance. The light vector mesons are introduced by the method of the hidden local symmetry. The weak current in the effective theory is obtained in the same way. The decay width of the process $D^{*+} \to D^0 \pi^+$, the form factors of the processes $\bar{B}^0 \to \pi^+ l \bar{\nu}$ and $\bar{B}^0 \to \rho^+ l \bar{\nu}$ are calculated. The results are easily translated to the D-meson semileptonic decays. The meson effective Lagrangian can give the simple and physically clear understanding on the heavy meson decays.

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FIGURES

FIG. 1. The diagrams for $\bar{B}^0 \to \pi^+ l \bar{\nu}$ decay. The black circle and box represent the strong and weak vertices, respectively.

FIG. 2. The diagrams for $\bar{B}^0 \to \rho^+ l \bar{\nu}$ decay.

REFERENCES

[1] N.Isgur and M.B.Wise, Phys. Lett. B232 (1989) 113; B237 (1990) 527.

[2] H.Georgi, Phys. Lett. B240 (1990) 447; M.J.Dugan, M.Golden, and B.Grinstein, Phys. Lett. B282 (1992) 142.

[3] H.Georgi, Lectures delivered at TASI, Published in Boulder TASI 91, 589 (HUTP-91-A039); G.Burdman and J.F.Donoghue, Phys. Lett. B280 (1992) 287; M.B.Wise, Phys. Rev. D45 (1992) R2188.

[4] M.E.Luke, Phys. Lett. B252 (1990) 447; M.Neubert, Phys. Lett. B306 (1993) 357; SLAC-PUB-6258.

[5] M.Luke and A.V.Manohar, Phys. Lett. B286 (1992) 348; M.Neubert, Phys. Lett. B306 (1993) 357.

[6] R.Casalbuoni, A.Deomdrea, N.Di Bartolomeo, R.Gatto, F.Feruglio, and G.Nardulli, Phys. Lett. B294 (1992) 106; R.Casalbuoni, N.Di Bartolomeo, R.Gatto, F.Feruglio, and G.Nardulli, Phys. Lett. B299 (1993) 139.

[7] T-M.Yan, H-Y.Cheng, C-Y.Cheng, G-L.Lin, Y.C.Lin, and H-L.Yu, Phys. Rev. D46 (1992) 1148.

[8] R.Casalbuoni, A.Deomdrea, N.Di Bartolomeo, R.Gatto, F.Feruglio, and G.Nardulli, Phys. Lett. B292 (1992) 371.

[9] M.Bando, T.Kugo, S.Uehara, K.Yamawaki, and T.Yanagida, Phys. Rev. Lett. 54 (1985) 1215; For review, M.Bando, T.Kugo, and K.Yamawaki, Phys. Rep. 164 (1988) 217.

[10] H-T.Cheng, C-Y.Cheung, G-L.Lin, Y.C.Lin, T-M. Yan, and H-L. Yu, CLNS 93/1192.

[11] CLEO collabolation, Phys. Rev. Lett. 69 (1992) 2041.

[12] ACCMOR collabolation, Phys. Lett. B278 (1992) 480.

[13] M.Neubert, SLAC-PUB-6344; G.Burdman, Z.Ligeti, M.Neubert, and Y.Nir, SLAC-PUB-6345.